

The Higgs as a Probe of Supersymmetric Extra Sectors

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Abstract

We present a general method for calculating the leading contributions to $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ in models where the Higgs weakly mixes with a nearly supersymmetric extra sector. Such mixing terms can play an important role in raising the Higgs mass relative to the value expected in the MSSM. Our method applies even when the extra sector is strongly coupled, and moreover does not require a microscopic Lagrangian description. Using constraints from holomorphy we fix the leading parametric form of the contributions to these Higgs processes, including the Higgs mixing angle dependence, up to an overall coefficient. Moreover, when the Higgs is the sole source of mass for a superconformal sector, we show that even this coefficient is often calculable. For appropriate mixing angles, the contribution of the extra states to $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ can vanish. We also discuss how current experimental limits already lead to non-trivial constraints on such models. Finally, we provide examples of extra sectors which satisfy the requirements necessary to use the holomorphic approximation.

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1 Introduction

The Higgs boson plays a privileged role in modern theories of particle physics, both as the last outstanding element of the Standard Model, and as a beacon for possible physics beyond the Standard Model (BSM). An attractive BSM scenario with both top-down and bottom-up motivations is supersymmetry. In particular, TeV scale supersymmetry provides an attractive way to stabilize the weak scale relative to the Planck scale.

In the absence of direct signatures of new physics at the weak scale, indirect signatures become all the more important. As has been appreciated for some time, the phenomenology of the Higgs sector itself provides a window into BSM physics. Indeed, processes such as $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ are generated by loop corrections, and thus are sensitive to heavy states which couple to both the Higgs and the massless $SU(3)_C \times U(1)_{EM}$ gauge bosons. Such effects are similar in spirit to other precision tests of the Standard Model.

In light of the above, the recent hints of an SM-like Higgs signal around 125 GeV by ATLAS [1] and CMS [2] are extremely exciting. For BSM scenarios such as the MSSM,

however, this leads to some tension with notions of naturalness since the tree level contribution to the Higgs quartic coupling arises from the $SU(2)_L$ and $U(1)_Y$ gauge couplings and is rather small. If the signal is real, getting a sufficiently heavy Higgs in the MSSM requires either large A-terms and/or heavy scalar superpartners (stops in particular) to raise the Higgs mass via radiative corrections. An alternative is to go beyond the MSSM, and consider setups with additional states which can provide further tree level and radiative contributions to the Higgs quartic potential. In addition to raising the Higgs mass to the observed level, these states can have other effects on Higgs physics, such as altering the loop level processes $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$. See [3–5] for some studies of the Higgs sector in scenarios beyond the MSSM.

With the above motivation in mind, we consider scenarios where the usual supersymmetric Higgs sector comprised of chiral superfields H_u and H_d mixes with a nearly supersymmetric extra sector via F- and D-terms. For example, the leading superpotential terms which can mix the two sectors are:

$$W_{mix} = \lambda_u H_u \mathcal{O}_u + \lambda_d H_d \mathcal{O}_d + \text{quadratic in } H\text{'s}. \quad (1.1)$$

where \mathcal{O}_u and \mathcal{O}_d are operators in some additional sector. Scenarios of electroweak symmetry breaking based on such mixing terms have been considered for example in [6–10]. Weakly coupled analogues involve adding a vector-like generation [11, 12]. More generally, the dynamics from an extra sector can introduce large additional corrections to the Higgs potential, which in particular can produce a much wider range of possible Higgs masses and mixing angles as compared with the MSSM. Examples include the Fat Higgs scenario [13], λ -SUSY [14] and the DSSM [10]. Independent of naturalness considerations (though not incompatible with them), the presence of additional sectors is also a common theme in various top-down motivated constructions such as [15].

There is clearly a huge range of possible extra sectors, which can run the gamut from weakly coupled to strongly coupled examples. Such extra sectors can potentially produce spectacular, though model dependent, signatures at the LHC. In many cases of interest, the extra sector may possess extra colored states, which could be light (around the electroweak scale) but still naturally evade the present bounds. The phenomenology of such states has been discussed in [16] as well as [10]. An interesting feature of adding such states is that it is necessary to include additional operators which mix with e.g. the third generation of the Standard Model, so that they can eventually decay¹. The focus of this work is on the

¹This can occur through the F-term $\Psi_R^{(3)} \cdot \mathcal{O}_{\bar{R}}$ between a third generation chiral superfield $\Psi_R^{(3)}$ and an operator $\mathcal{O}_{\bar{R}}$ with conjugate quantum numbers. Fortunately, such couplings are automatically present in string constructions such as [15]. This may lead to the impression that if the spectrum of the extra sector comes in the form of full GUT multiplets to preserve gauge unification, then this could lead to fast proton decay via operators of the form $QQQL/M_{extra}$, generated for example by integrating out colored triplet states in the extra sector with masses around the TeV scale (if no symmetry suppresses it). However, it can

indirect effects of these states on Higgs physics. Indeed, the extra sector may be hard to probe directly, but could still have consequences for Higgs physics. Of course, the (model dependent) collider phenomenology of these states should be explored further in the future.

When the masses of extra states m_i are sufficiently heavy ($m_h^2 \ll 4m_i^2$), their effects on Higgs couplings can be included via higher dimension operators such as:

$$\mathcal{O}_{hFF} = c \cdot \frac{h^0}{v} \text{Tr}_G F^2 \quad (1.2)$$

where $G = SU(3)_C, U(1)_{EM}$, c is an “order one number”, and $v \sim 246$ GeV is the Higgs vev. It is well known that the general form of this contribution can be extracted from the gauge coupling threshold correction due to the extra states [17]. The detailed form of this threshold, however, depends on the mass spectrum of the extra states, and so can be difficult to extract in general.

In the limit where the extra sector is approximately supersymmetric, a great deal of information about \mathcal{O}_{hFF} and related dimension five operators can be extracted. In models which admit vector-like masses, we can consider adding gauge invariant mass deformations which decouple all of the extra sector states. Holomorphy and gauge invariance then dictate the form of the leading-order contribution to $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ from the dimension six operator:

$$\mathcal{L}_{eff} \supset \text{Re} \frac{-b_G}{8\pi^2} \int d^2\theta \frac{H_u H_d}{\Lambda_G^2} \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha \quad (1.3)$$

where Λ_G is a characteristic mass scale, and b_G is a dimensionless constant we shall identify with the beta function coefficient contribution from the extra states. This leads to the dimension five operator

$$\mathcal{O}_{hFF} = \frac{b_G}{16\pi^2} \cdot \cos(\alpha + \beta) \cdot \left(\frac{v}{\Lambda_G}\right)^2 \cdot \frac{h^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} \quad (1.4)$$

where α and β are the Higgs mixing angles with conventions as in [18].

In fact, in many cases even more is known about the form of this dimension five operator. For example, when the extra sector is a superconformal field theory (SCFT), b_G is often a *calculable* global anomaly coefficient; we review this fact in section 2. Moreover, when the

be shown that the coefficients of these operators are sufficiently suppressed in many interesting cases. For example in a superconformal extra sector generating $QQQL/M_{extra}$ involves correlators of at least four \mathcal{O} operators. Setting M_{GUT} as the UV cutoff scale and M_{extra} as the IR cutoff there is order $(M_{extra}/M_{GUT})^D$ conformal suppression, where $D \sim 4 \times 2$ for operators \mathcal{O} of dimension close to two. Such contributions are below all conceivable detection limits.

Higgs is the sole source of mass, we have:

$$\mathcal{O}_{hFF} = \frac{1}{16\pi^2} \left(b_u \frac{\cos \alpha}{\sin \beta} - b_d \frac{\sin \alpha}{\cos \beta} \right) \cdot \frac{h^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} \quad (1.5)$$

where b_u and b_d are again threshold coefficients, which are often calculable when the extra sector is superconformal. In many well-motivated situations, $b_u = b_d = b_G/2$, which reduces to equation (1.4) when $\Lambda_G^2 = 2v_u v_d$.

Aside from being a particularly calculable limit, the case of superconformal extra sectors is also attractive because it can allow for large Yukawa couplings without the worry of a low Landau pole (as the running stops once we enter the conformal regime). Further, for appropriate CFTs, it is possible to have large Higgs-extra sector Yukawas while still maintaining small anomalous dimensions for the Higgs fields, a point we discuss further in section 4.

When applicable, the holomorphic approximation clearly provides a powerful constraint on the possible contributions of extra sectors to Higgs physics. One of our tasks in this paper will be to estimate the expected regime of validity; subleading corrections can occur in both the supersymmetric limit as well as in the limit where supersymmetry is broken. We find that the main criterion which must be satisfied is that the anomalous dimension of the Higgs must be small. In this limit, the Higgs retains its identity as a weakly coupled field.² Fortunately, this is also the regime which is favored by current limits on extensions of the Standard Model. Further, in this regime perturbative visible sector gauge and Yukawa couplings can be maintained.

The rest of this paper is organized as follows. In section 2 we present the basic idea of the holomorphic approximation, and detail the expected regime of validity. Next, in section 3 we compare with experiment, illustrating the utility of the method as a constraint on possible Higgs-extra sector mixing. In section 4 we provide some explicit examples of supersymmetric extra sectors which satisfy the criteria necessary to use the holomorphic approximation. In particular, we find that scenarios inspired by F-theory are a particularly attractive class of models. We present our conclusions in section 5.

2 The Holomorphic Approximation

In this section we explain how to extract the leading-order dimension five operators from F-term data. We refer to this as the *holomorphic approximation*, since the dominant couplings will be extracted from holomorphic data.

²This situation should be contrasted to one in which the Higgs picks up a large anomalous dimension and thus is better viewed as a composite.

Our basic setup is as follows. We view the Standard Model gauge group as a flavor symmetry of an extra sector which may exhibit strong coupling dynamics. We assume, however, that the mass spectrum in the extra sector is approximately supersymmetric. Additionally, we wish to remain in a regime where to leading order the Higgs vevs can be treated as spurions. In this limit, we can track how the weakly gauged SM “flavor symmetries” $SU(3)_C \times U(1)_{EM}$ respond to the Higgs vevs. Using this information, we will extract the leading-order contribution to the dimension five operator $h^0 \text{Tr}_G F^2$.

It is well known that in the limit where the masses m_i of these states are large compared to the Higgs mass ($m_h^2 \ll 4m_i^2$), the contribution from the extra states to the dimension five operator $h^0 \text{Tr}_G F^2$ can be modelled as a threshold correction to the $SU(3)_C$ and $U(1)_{EM}$ gauge couplings [17]:

$$\mathcal{O}_{hFF} = \frac{b_G}{32\pi^2} \left(v \frac{\partial \log \det \mathcal{M}}{\partial v} \right) \frac{h^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu}, \quad (2.1)$$

where b_G is the beta function coefficient from the threshold characterized by \mathcal{M} , the mass matrix of the extra states. In a two-Higgs doublet model (2HDM), the mass matrix \mathcal{M} can depend on the vevs v_u and v_d in a complicated way. This is especially true for a strongly coupled extra sector, where little quantitative information is typically available. It would be useful to learn about how the extra sector fixes $h^0 \text{Tr}_G F^2$ *without* a detailed analysis of the extra sector mass spectrum and couplings, as they will be difficult to measure (especially at hadron colliders).

When the extra sector is nearly supersymmetric and mixes weakly with the Higgs sector, additional constraints come into play. Our main focus will be on the limit where we add vector-like $SU(3)_C \times SU(2)_L \times U(1)_Y$ preserving mass terms to the extra sector. In this case, we can integrate out these states, to generate the dimension six F-term:

$$\mathcal{L}_{eff} \supset \text{Re} \frac{-b_G}{8\pi^2} \int d^2\theta \frac{H_u H_d}{\Lambda_G^2} \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha \quad (2.2)$$

where Λ_G is a characteristic mass scale, and b_G is a dimensionless constant we shall identify with the beta function coefficient contribution from the extra states (see subsection 2.1). Here, the gauge kinetic term is given by:

$$\mathcal{L}_{kin} = \text{Im} \frac{\tau}{8\pi} \int d^2\theta \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha = -\frac{1}{2g^2} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} + \frac{\theta}{32\pi^2} \varepsilon^{\mu\nu\rho\sigma} \text{Tr}_G F_{\mu\nu} F_{\rho\sigma} \quad (2.3)$$

where $\tau = \frac{4\pi i}{g^2} + \frac{\theta}{2\pi}$ is the holomorphic gauge coupling. In the limit where the Higgs-extra sector Yukawas can be treated as perturbative, we have the further relation:

$$M_{extra}^2 \sim \lambda_u \lambda_d \Lambda_G^2 \quad (2.4)$$

where M_{extra} are the masses of the extra sector states.

Quite remarkably, this is enough to fully fix the Higgs mixing angle dependence. Indeed, expanding in the mass eigenstate basis:³

$$h_u^0 = \frac{1}{\sqrt{2}} (v_u + h^0 \cos \alpha + H^0 \sin \alpha + iA^0 \cos \beta + \text{Goldstones}) \quad (2.5)$$

$$h_d^0 = \frac{1}{\sqrt{2}} (v_d - h^0 \sin \alpha + H^0 \cos \alpha + iA^0 \sin \beta + \text{Goldstones}) \quad (2.6)$$

we obtain a remarkably rigid expression for the form of the dimension five operators:

$$\mathcal{O}_{hFF} = \frac{b_G}{16\pi^2} \cdot \cos(\alpha + \beta) \cdot \left(\frac{v}{\Lambda_G}\right)^2 \cdot \frac{h^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} \quad (2.7)$$

$$\mathcal{O}_{HFF} = \frac{b_G}{16\pi^2} \cdot \sin(\alpha + \beta) \cdot \left(\frac{v}{\Lambda_G}\right)^2 \cdot \frac{H^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} \quad (2.8)$$

$$\mathcal{O}_{AFF} = \frac{b_G}{32\pi^2} \cdot \left(\frac{v}{\Lambda_G}\right)^2 \cdot \frac{A^0}{v} \varepsilon^{\mu\nu\rho\sigma} \text{Tr}_G F_{\mu\nu} F_{\rho\sigma}. \quad (2.9)$$

Observe also that the contributions decouple as $(v/\Lambda_G)^2$ since they descend from a supersymmetric dimension six operator. Note also that the ratios of the couplings for the CP-even and odd states are all completely fixed, depending only on the Higgs mixing angles.

Clearly, when it applies, the holomorphic approximation leads to a remarkably rigid structure on the possible contributions to the Higgs sector. In the remainder of this section, we explain how this approximation can be viewed as a supersymmetric threshold, and moreover, how to calculate the coefficient b_G . After this, we show that the *exact* form of \mathcal{O}_{hFF} can be extracted when the Higgs is the sole source of mass for a superconformal extra sector. Finally, we discuss the expected regime of validity in the presence of supersymmetry breaking.

2.1 Supersymmetric Thresholds

In this subsection we show that equation (2.2) originates as a supersymmetric threshold correction to the visible sector gauge couplings. Our main assumption here will be that the mass spectrum of the extra states is nearly supersymmetric. Further, we assume that the extra sector admits vector-like mass deformations, i.e. mass terms which preserve $SU(3)_C \times SU(2)_L \times U(1)_Y$.

Assuming supersymmetry is not broken, it is convenient to work in a formalism where

³Throughout this paper, we assume that in the Higgs sector, CP is conserved so that h^0 , H^0 are CP-even, and A^0 is CP-odd. This is reflected in v_u and v_d being real, and also feeds into the assumption that $\Lambda_G > 0$. In our conventions $v_u = v \sin \beta$, $v_d = v \cos \beta$, with $v = 246$ GeV.

all couplings and masses are treated as superfields. The point is that for unbroken gauge symmetry generators, holomorphy imposes a strong constraint on the possible couplings one can write. Promoting the holomorphic gauge coupling to a chiral superfield yields the F-term coupling:

$$\mathcal{L}_{\tau\mathcal{W}\mathcal{W}} = \text{Im} \int d^2\theta \frac{\tau(\mu)}{8\pi} \cdot \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha. \quad (2.10)$$

It is well-known that in the holomorphic basis of fields, τ is exact at at one loop and satisfies:

$$\tau = \tau_0 + \frac{b_G^{(h)}}{4\pi i} \log \frac{M^2}{\mu_0^2} \quad (2.11)$$

where τ_0 is the value of τ at the reference scale μ_0 , M corresponds to a mass threshold, and $b_G^{(h)}$ is the holomorphic beta function coefficient corresponding to the supersymmetric mass threshold M . The general form of these couplings will then be specified in terms of a holomorphic function $M^2(X_H, X_i)$ with $X_H \equiv H_u H_d$.⁴

$$\mathcal{L}_{eff} \supset \text{Re} \int d^2\theta \frac{-b_G}{32\pi^2} \log M^2(X_H, X_i) \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha. \quad (2.12)$$

Thematically, this is similar to the idea of analytic continuation in superspace often employed in minimal gauge mediation [19, 20]. To read off the leading-order couplings to the Higgs fields in this limit, we expand $M^2(X_H, X_i)$ to linear order in the Higgs field vevs:

$$\mathcal{L}_{eff} \supset \text{Re} \int d^2\theta \frac{-b_G}{8\pi^2} \left(\frac{h_u^0 v_d}{\sqrt{2}\Lambda_G^2} + \frac{h_d^0 v_u}{\sqrt{2}\Lambda_G^2} \right) \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha. \quad (2.13)$$

Here, we have absorbed the Higgs-extra sector Yukawas into the definition of the characteristic scale Λ_G to retain the interpretation of b_G as a beta function coefficient. Expanding in the mass eigenstate basis, we recover equations (2.7)-(2.9). Let us note in passing that one can also expand in the moduli X_i to extract the leading-order X_i - F^2 couplings. For related discussion of pseudo-dilaton- F^2 couplings, see for example [21].

Even in the supersymmetric limit there can be additional non-holomorphic dependence on the Higgs fields. Indeed, to get the *physical* hFF vertex we must pass to a basis of canonically normalized superfields. We refer to \mathcal{W}^α as the gauge field strength in a holomorphic basis of fields, and by contrast, we reserve W^α for the gauge field strength in the “physical” i.e. canonically normalized basis of fields. The reason for this distinction is that when we go to the canonical basis of fields, the overall normalization will generically involve an anomalous non-holomorphic rescaling of W^α . The holomorphic approximation

⁴When H_u and H_d mix with a CFT and obtain dimensions Δ_{H_u} and Δ_{H_d} , we would make the replacement $\mu_0^2 \rightarrow \mu_0^{2\Delta}$ for some $\Delta > 1$. This effect is absorbed into the definition of the beta function coefficient b_G . See e.g. [15] for further discussion.

is a good one precisely when this subtlety can be ignored.

Such effects are encapsulated in the more general expression which contains the gauge kinetic term (see e.g. [22] for discussion of this term in the context of “gaugino screening”):

$$\mathcal{L}^{(c)} \supset \int d^4\theta \frac{\Omega(\mu)}{8\pi} \text{Tr}_G W^\alpha \left(\frac{D^2}{-8p^2} \right) W_\alpha \quad (2.14)$$

where $\Omega(\mu)$ is a real superfield related to τ via:

$$\Omega(\mu) = \text{Im } \tau(\mu) - \frac{1}{2\pi} \sum_i t_2^i \log \mathcal{Z}_i(\mu) + \dots \quad (2.15)$$

Here, $\mathcal{Z}_i(\mu)$ is the contribution from wavefunction renormalization and i runs over the Higgs fields and all states charged under the visible gauge couplings ($SU(3)$ or $U(1)_Y$). The “...” are terms involving the gauge coupling of G , and are suppressed because the visible gauge couplings are perturbative. Whereas the τ -dependent terms are manifestly holomorphic, the contributions $\mathcal{Z}_i(\mu)$ include all of the non-holomorphic contributions to the masses. In $\Omega(\mu)$, the contribution $t_2^i \log \mathcal{Z}_i(\mu)$ is summed over the matter fields, and assuming a threshold scale m_i for each species can be written as:

$$\sum_i t_2^i \log \mathcal{Z}_i(\mu) \simeq \sum_i t_2^i \gamma^i \log |m_i| \simeq \delta b_G \log \det \mathcal{M} \quad (2.16)$$

where \mathcal{M} is the mass matrix for the extra states, and we have introduced $\delta b_G \equiv b_G^{(NSVZ)} - b_G^{(h)}$, the difference between the NSVZ beta function and the beta function of the holomorphic gauge coupling. Indeed, equation (2.15) contains the same physical content as the numerator of the famous NSVZ beta function [23, 24], which is also sensitive to non-holomorphic contributions through the anomalous dimensions of the fields. If $\delta b_G/b_G$ is small, then one can expand in this parameter. Making the formal replacement $\mathcal{W}^\alpha \rightarrow W^\alpha$, the size of the correction term in equations (2.7)-(2.9) will be of order $\delta b_G/b_G$ relative to the term multiplying b_G .

2.2 The Coefficient b_G

As we have seen, the holomorphic approximation is helpful precisely when $\delta b_G/b_G \ll 1$. At an intuitive level, the ratio $\delta b_G/b_G$ quantifies the amount of Higgs-extra sector mixing. In this subsection we make this intuition more precise, and explain why the small mixing regime is phenomenologically favored. Additionally, we explain at an abstract level how to compute both b_G and $\delta b_G/b_G$ in the special case where the extra sector is superconformal. In particular, this means that even in the strongly coupled setting, it is possible to compute

the “order one coefficient” multiplying $h^0 \text{Tr}_G F^2$.

Let us begin by quantifying the amount of mixing between the Higgs and the extra sector. For our purposes, this is captured by the shift in the anomalous dimension of the Higgs fields, as well as the operators of the extra sector. Since these anomalous dimensions also show up in the numerator of the NSVZ beta function, we can track the amount of mixing through changes to the beta function. To this end, we consider three theories associated with our extra sector. As usual, we work in the approximation where all Standard Model fields (except the Higgs fields) are treated as non-dynamical. One theory is given by a “UV theory”, in which all couplings to the Higgs have been switched off. We also consider a “Mixed theory” in which the couplings between the Higgs and the extra sector have been switched on. Finally, we consider an “IR theory” in which we have activated a Higgs vev. For each of these theories, we can weakly gauge our flavor symmetry group, and compute the resulting beta function coefficient at a scale μ . We say there is little mixing between the two sectors when $\delta b = b^{UV} - b^{MIX}$ is small compared to b^{UV} and b^{MIX} . We also see that the size of the threshold correction b_G is given by $b^{MIX} - b^{IR}$, as this corresponds to the threshold correction from all states which can get a mass from the Higgs coupling. Note that in many situations of interest where the Higgs vev gives a mass to all states, $b^{IR} = 0$. On the other hand, one can also contemplate scenarios where only some of the states of the extra sector directly couple to the Higgs. The difference $b^{MIX} - b^{IR}$ quantifies this contribution.

It is important to distinguish here between the mixing induced by the beta functions, $\delta b_G/b_G$, and that associated with Yukawa couplings such as $\lambda_u H_u \mathcal{O}_u + \lambda_d H_d \mathcal{O}_d$. This is because one can consider situations where the Higgs develops only a small anomalous dimension even though λ_u and λ_d may be large. We will discuss examples of this type in section 4.

In actual applications, we are interested in the value of the beta function coefficients for $G = SU(3)_C$ and $U(1)_{EM}$. In terms of the beta functions for $SU(2)_L$ and $U(1)_Y$, we have:

$$b_{EM} = b_{SU(2)} + \frac{5}{3} b_{U(1)} \quad (2.17)$$

where we have normalized $U(1)_Y$ so that it is embedded in $SU(5)_{GUT}$. Observe that in the special case where b^{UV} retains gauge coupling unification, we have $b_{GUT} = b_{SU(3)} = b_{SU(2)} = b_{U(1)}$, so that $b_{EM} = \frac{8}{3} b_{GUT}$. Away from the vector-like mass limit, one can also track the dependence on just v_u and just v_d , and two corresponding threshold coefficients b_u and b_d .

A fortunate feature of the holomorphic approximation is that it works best in the limit of small mixing where $\delta b_G/b_G \ll 1$, which is also the regime favored by various phenomenological considerations. Indeed, the larger the mixing between the Higgs and the extra sector, the more the Higgs will deviate from a weakly coupled scalar. This is disfavored by various

indirect precision electroweak measurements, as well as by the (still accumulating) evidence for a relatively light Higgs boson. Additionally, when the Higgs field has dimension greater than one, maintaining relatively large Yukawa couplings with other Standard Model fields becomes more tenuous. Conversely, when the Higgs has dimension less than one (as could happen in a CFT), this requires $SU(2)_L \times U(1)_Y$ to become strongly coupled.⁵ Maintaining small $\delta b_G/b_G$ can also help with gauge coupling unification. This is because the Higgs fields do not fill out complete GUT multiplets, so that large mixing could distort gauge coupling unification. See e.g. [25] for further discussion on this point.

Finally, when the extra sector is a superconformal field theory, it is often possible to *calculate* b_G , even without a Lagrangian description of the extra sector. This is because b_G is actually a global anomaly coefficient:

$$b_G \delta^{ab} = -3 \text{Tr}(R_{IR} J_G^a J_G^b), \quad (2.18)$$

where R_{IR} is the superconformal R-current and J_G^a is a global symmetry current which we weakly couple to the standard model vector multiplet V_{SM} by $\mathcal{L} = \int d^4\theta V_{SM} \mathcal{J}_G$ (\mathcal{J}_G is the current superfield containing J_G , see e.g. [26]). The key point is that this beta function coefficient can often be computed via 't Hooft anomaly matching, as in [25]. In section 4 we provide some further examples where we calculate such contributions. Note also that when the extra sector is an SCFT, an important and model-independent unitarity constraint is that $b_G > 0$ (see e.g. [27, 28]).

2.3 Away from the Vector-Like Limit

In this subsection, we consider situations where the extra sector states may not possess large vector-like masses. Perhaps surprisingly, in the limit where the Higgs is the sole source of mass for a superconformal extra sector we show that the *exact* form of the dimension five operator is fixed by visible sector parameters.

At the level of the effective field theory, the most general dimension five operator consistent with holomorphy is:

$$\mathcal{L}_{int} = \text{Re} \int d^2\theta \left(-\frac{b_u}{16\pi^2} \frac{h_u^0}{\Lambda_u} - \frac{b_d}{16\pi^2} \frac{h_d^0}{\Lambda_d} \right) \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha \quad (2.19)$$

where Λ_u and Λ_d are characteristic mass scales, and b_u and b_d are the beta function coefficients for states which get their mass from v_u and v_d , respectively. Here, as earlier, we

⁵Indeed, we would arrive at contradiction if we allowed $SU(2)_L \times U(1)_Y$ to remain as a weakly gauged flavor symmetry with H_u or H_d becoming a gauge-invariant operator with dimension below the unitarity bound. Hence, it is necessary to allow the weakly gauged flavor symmetry to instead become strongly coupled. This is a logical, though unappealing possibility.

implicitly assume that the Higgs sector preserves CP, so that we can take b_u/Λ_u , b_d/Λ_d , v_u and v_d all real. Expanding in the mass eigenstate basis we can read off the couplings to all three electrically neutral states. In contrast to a general two Higgs doublet model, holomorphy allows us to relate the dimension five operators for all three electrically neutral states in terms of two undetermined coefficients, b_u/Λ_u and b_d/Λ_d .

The situation becomes far more predictive when the Higgs fields are the sole source of mass for states of a superconformal extra sector. Although this limit does lead to some tension with constraints from precision electroweak data, viable scenarios exist which satisfy all current bounds [10]. Our main interest in this case here is that it is a remarkably calculable limit. Indeed, the form of the supersymmetric threshold in this special case is:

$$\mathcal{L}_{int} = \text{Re} \int d^2\theta \left(-\frac{b_u}{16\pi^2} \log \frac{h_u^0}{\mu_0} - \frac{b_d}{16\pi^2} \log \frac{h_d^0}{\mu_0} \right) \text{Tr}_G \mathcal{W}^\alpha \mathcal{W}_\alpha. \quad (2.20)$$

For a superconformal extra sector, the coefficients b_u and b_d are specified as follows. Once we switch on either v_u or v_d , we introduce a relevant deformation of the theory. This leads to a new IR theory, with corresponding beta functions b_u^{IR} and b_d^{IR} for the two cases. We identify $b_u = b^{MIX} - b_u^{IR}$ and $b_d = b^{MIX} - b_d^{IR}$ (as in subsection 2.2). Expanding in the Higgs mass eigenstates, we obtain the explicit form of the dimension five operators:

$$\mathcal{O}_{hFF} = \frac{1}{16\pi^2} \left(b_u \frac{\cos \alpha}{\sin \beta} - b_d \frac{\sin \alpha}{\cos \beta} \right) \cdot \frac{h^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} \quad (2.21)$$

$$\mathcal{O}_{HFF} = \frac{1}{16\pi^2} \left(b_u \frac{\sin \alpha}{\sin \beta} + b_d \frac{\cos \alpha}{\cos \beta} \right) \cdot \frac{H^0}{v} \text{Tr}_G F_{\mu\nu} F^{\mu\nu} \quad (2.22)$$

$$\mathcal{O}_{AFF} = \frac{1}{32\pi^2} \cdot (b_u \cot \beta + b_d \tan \beta) \cdot \frac{A^0}{v} \varepsilon^{\mu\nu\rho\sigma} \text{Tr}_G F_{\mu\nu} F_{\rho\sigma}. \quad (2.23)$$

All dependence on the Higgs-extra sector Yukawas has dropped out. Indeed, everything has reduced to a computation of the *calculable* coefficients b_u and b_d . In the special – though well-motivated – case where $b_u = b_d = b_G/2$, observe that the parametric form collapses further to equations (2.7)-(2.9) with the replacement:

$$\Lambda_G^2 = 2v_u v_d = v^2 \sin 2\beta. \quad (2.24)$$

2.4 Incorporating Supersymmetry Breaking

So far, we have worked in a limit where the extra sector is supersymmetric. In the more realistic case, there will be supersymmetry breaking contributions to the masses, encapsulated in F-term components of X_H and X_i , with notation as in subsection 2.1. For our approximation to be valid, the F-term components of these spurions must be a subleading

contribution, relative to the square of their scalar components.

One source of supersymmetry breaking is unavoidable, coming from the Higgs F -term vevs:

$$\langle h_u^0 \rangle = \frac{1}{\sqrt{2}} (v_u + \theta^2 F_u), \quad \langle h_d^0 \rangle = \frac{1}{\sqrt{2}} (v_d + \theta^2 F_d) \quad (2.25)$$

in the obvious notation. This feeds into the extra sector through F -term couplings such as:

$$\mathcal{L}_{mix} = \int d^2\theta (\lambda_u H_u \mathcal{O}_u + \lambda_d H_d \mathcal{O}_d) + h.c.. \quad (2.26)$$

The supersymmetry breaking contributions will be small provided:

$$(\lambda_u v_u)^2 \gg \lambda_u F_u, \quad (\lambda_d v_d)^2 \gg \lambda_d F_d. \quad (2.27)$$

where $M_u = \lambda_u v_u$ and $M_d = \lambda_d v_d$ are the characteristic mass scales of states of the extra sector.⁶ This is similar to the case of messengers in gauge mediation with the Higgs replaced by a SUSY breaking spurion.

Using $F_u \sim \mu v_d$, $F_d \sim \mu v_u$ where μ is the supersymmetric mass term of the Higgs sector, these conditions become:

$$M_u^2 \tan^2 \beta \gg \mu^2, \quad M_d^2 \cot^2 \beta \gg \mu^2. \quad (2.28)$$

For both up-type and down-type states to be sufficiently heavy, a natural possibility is $M_u \sim M_d$ and $\tan \beta \sim O(1)$ (although some hierarchy between M_u and M_d as well as correspondingly large $\tan \beta$ are also possible). As an example, taking $M_u \sim M_d \sim 1$ TeV and $\mu \sim 200$ GeV with $\tan \beta = 1$, we have $(\mu/M_u)^2 \sim 0.04$.

Consider next supersymmetry breaking contributions from sources other than the Higgs. Here, the situation is clearly more model dependent. However, we find that supersymmetry breaking mediation mechanisms are often compatible with having a nearly supersymmetric extra sector. To illustrate the point, consider the case where the extra sector is approximately conformal, but the visible sector has superpartner masses on the order of ~ 1 TeV.⁷ We would like to know how supersymmetry breaking will be transmitted to the states of

⁶As noted in [10], at strong coupling we do not really know the mass of the extra states. However, it is reasonable to expect that they are proportional to the Yukawa coupling of the hidden sector. At weak coupling, the mass of the extra states is proportional to $\sqrt{\delta}$, where δ is the excess Higgs dimension. This provides a conservative (though rough) lower bound on the mass of such extra states.

⁷Let us note that one can still contemplate rather light visible sector superpartners, which may have evaded detection thus far. In such cases, the holomorphic approximation applies if the extra sector states have TeV scale masses (as can happen from having large Higgs-extra sector Yukawas). Examples include compressed superpartner spectra or R-parity violating models. Additionally, in string constructions of extra sectors such as [15], there can in principle be geometric sequestering between the location of the extra sector and a possibly localized supersymmetry breaking sector.

the extra sector, and in particular, whether the dominant contribution to the masses of states will be from supersymmetry preserving terms such as the vector-like mass and Higgs vevs (for sufficiently large Higgs-extra sector Yukawas) or will instead be dominated by supersymmetry breaking effects. An important point to keep in mind is that the Green’s function for a state of the extra sector will typically deviate from the free field expression, being instead given by an “unparticle propagator” (see e.g. [29]) which for a scalar operator takes the form:

$$\langle \mathcal{U}^\dagger(x) \mathcal{U}(0) \rangle \sim \frac{1}{(x^2)^\Delta}. \quad (2.29)$$

Unitarity requires $\Delta > 1$ (see e.g. [30, 31]). These Green’s functions feed into the transmission of supersymmetry breaking in the extra sector. In particular, relative to the soft mass scale Λ_{soft} of visible sector states, there will be additional suppression factors of order $(M_{CFT}/M_{mess})^{\Delta-1}$ for the soft masses of the extra sector, where M_{CFT} is the CFT breaking scale. This is of course a well known phenomenon in the context of conformal sequestering (see e.g. [32–34]), though here the motivation and application of this phenomenon is somewhat different. To give a numerical example, consider $M_{CFT} \sim 1$ TeV and $\Delta \sim 1.1$. This yields a factor of ten suppression in the extra sector supersymmetry breaking mass terms when $M_{mess} \sim 10^{13}$ GeV, as can happen in intermediate scale gauge mediation models. This suffices for the supersymmetric mass terms to dominate, and illustrates that the extra sector can naturally shield itself from supersymmetry breaking effects, so that the holomorphic approximation applies.

3 Higgs Phenomenology

The recent hints of a Standard Model-like Higgs with a mass close to 125 GeV are very exciting. Assuming that the signal is real and is due to an h^0 resonance, it is of crucial theoretical interest to figure out if data in the various channels measured by ATLAS and CMS could be used as a probe of BSM physics. In this section we study this issue for a Higgs which couples to a supersymmetric extra sector. Further, we work under the assumption that there is a vector-like mass in the extra sector, so that we can potentially decouple the presence of such states (though we do not work in that limit). In this case, the parametric form of equations (2.7)-(2.9) applies. It is simple to also interpret our results in the case where the Higgs is the sole source of mass for a vector-like conformal sector using the substitution (2.24). However, one should keep the following caveats in mind when interpreting our results:

- The present data on various channels is rather preliminary and could change significantly, both in terms of central values and/or uncertainties. This could happen due to an upward fluctuation in signal (which is not uncommon when looking for

a new signal), or due to an improvement in understanding systematic uncertainties. An interesting example is $pp \rightarrow h^0 jj \rightarrow \gamma\gamma jj$. After imposing relevant cuts, this channel gets a large contribution from vector boson fusion (VBF). However, gluon fusion with two radiated jets also provides an important contribution which is not precisely known and could have significant uncertainties. We expect that more data will improve the situation considerably.

- We only focus on search channels associated with the Higgs signal, and do not perform an analysis of other LHC searches for the colored and electroweak states in the extra sector, since signatures of such states are quite model dependent. Indeed, part of our point is that even without knowing all of these details, the Higgs itself is an excellent probe of such sectors.
- The amplitudes for the processes $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ can be viewed as the sum of three contributions which, normalized relative to the Standard Model, can be written as:

$$\frac{A}{A_{SM}} = \hat{A}_{s2HDM} + \hat{A}_{MSSM} + \hat{A}_{Extra} \quad (3.1)$$

where \hat{A} is the ratio of amplitudes A/A_{SM} . \hat{A}_{s2HDM} denotes the contributions from the supersymmetric 2HDM, \hat{A}_{MSSM} denotes the contribution from all superpartners in the MSSM (or an extension thereof), and \hat{A}_{Extra} denotes a possible contribution from the extra states, all normalized relative to the SM contribution.

The contribution from \hat{A}_{MSSM} decouples as v^2/M_{SUSY}^2 for soft masses $M_{SUSY}^2 > v^2$, while that from \hat{A}_{Extra} decouples as v^2/Λ^2 . In the vector-like mass limit the parametric dependence on the Higgs angles is fixed, and further suppression occurs in the limit $\cos(\alpha + \beta) \rightarrow 0$ (see (1.4)).

In this work, for simplicity, we study the case where the contribution \hat{A}_{MSSM} is decoupled but \hat{A}_{Extra} is not, so that data can constrain the properties of the extra sector in a simple manner. The bounds on superpartners keep getting better, so our assumption may be well justified. However, it is worth noting that current data still allows comparatively light third generation squarks and sleptons which could have an important effect on $h^0 \rightarrow gg$ (see e.g. [35]) and $h^0 \rightarrow \gamma\gamma$ (see e.g. [36]) respectively.

3.1 Higgs Partial Widths

Before discussing implications of the data on our setup, it is useful to collect the relevant expressions for the various Higgs couplings relative to the SM, and set up the notation. We introduce quantities γ_{ii} defined as the h^0 partial width to the state ii normalized to the

SM Higgs partial width to the same final state:

$$\gamma_{ii} \equiv \frac{\Gamma(h^0 \rightarrow ii)}{\Gamma(h_{SM} \rightarrow ii)}. \quad (3.2)$$

The total width $\Gamma_{h^0}^{tot}$ and the cross-section for a given channel ($XX \rightarrow h^0 \rightarrow ii$) relative to the SM are given by:

$$\begin{aligned} R_\Gamma &\equiv \frac{\Gamma_{h^0}^{tot}}{\Gamma_{h_{SM}}^{tot}} = \frac{\sum_{ii} \Gamma(h^0 \rightarrow ii)}{\Gamma_{h_{SM}}^{tot}} = \sum_{ii} (B_{SM}^i \gamma_{ii}) \\ R_{Xi} &\equiv \frac{\sigma(XX \rightarrow h^0 \rightarrow ii)}{\sigma(XX \rightarrow h_{SM} \rightarrow ii)} = \frac{1}{R_\Gamma} \left(\frac{\gamma_{XX}}{\gamma_{ii}} \right). \end{aligned} \quad (3.3)$$

Here B_{SM}^i is the branching ratio of the SM Higgs to final state ii . The notation is similar to that of [37]. Note that in (3.3), we have assumed that the Higgs does not have an invisible decay width. Although in principle one can contemplate decays of the Higgs to hidden sector singlets or neutralino LSPs which increase the total Higgs width and in turn lower the various branching fractions, this generically lowers the expected signal (though it can compensate for an increase in a production channel).

In terms of the above quantities, the Higgs decay widths to up- and down-type fermions ($f_u \bar{f}_u$, $f_d \bar{f}_d$) and massive vector bosons ($VV = WW, ZZ$) are dominated by tree-level decays, and will be essentially the same as in the usual supersymmetric 2HDM (see [38–40] for reviews):⁸

$$\gamma_{f_u \bar{f}_u} = \left(\frac{\cos \alpha}{\sin \beta} \right)^2 ; \quad \gamma_{f_d \bar{f}_d} = \left(-\frac{\sin \alpha}{\cos \beta} \right)^2 ; \quad \gamma_{VV} = \sin^2(\beta - \alpha). \quad (3.4)$$

On the other hand, the loop processes $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ will be sensitive to the contributions from the extra states.⁹ In the SM, the dominant contributions to these processes are respectively from a top quark loop and a W -boson loop. The widths for these

⁸Let us note that the mixing with the extra sector can induce corrections to the Kähler potential for the Higgses. This shows up as a modification of the mass of the SM states as a function of v_u and v_d . However, in the limit where such wave function renormalization effects are small (which is necessary to apply the holomorphic approximation anyway), this shift is also small. See [10] for further discussion.

⁹This is also true for the process $h^0 \rightarrow Z\gamma$ which will soon play an important role in Higgs searches. A subtlety with applying the holomorphic approximation in this case is that the effective operator involves a gauge boson of a broken symmetry generator.

processes relative to the SM are to leading order given by:¹⁰

$$\gamma_{gg} \simeq \frac{\left| \frac{\cos \alpha}{\sin \beta} A_{1/2}(\tau_t) - \frac{\sin \alpha}{\cos \beta} A_{1/2}(\tau_b) + A_{Extra}^{gg} \right|^2}{|A_{1/2}(\tau_t)|^2}; \quad (3.5)$$

$$\gamma_{\gamma\gamma} \simeq \frac{\left| \sin(\beta - \alpha) A_1(\tau_W) + \frac{4}{3} \frac{\cos \alpha}{\sin \beta} A_{1/2}(\tau_t) - \frac{1}{3} \frac{\sin \alpha}{\cos \beta} A_{1/2}(\tau_b) - \frac{\sin \alpha}{\cos \beta} A_{1/2}(\tau_\tau) + A_{Extra}^{\gamma\gamma} \right|^2}{|A_1(\tau_W) + \frac{4}{3} A_{1/2}(\tau_t)|^2}. \quad (3.6)$$

The loop contribution from H^\pm is relatively small, so we do not include it in what follows. Here $\tau_i \equiv \frac{m_h^2}{4m_i^2}$, and $A_s(\tau_i)$ is a form factor for a particle in the loop with spin s and mass m_i [38, 39]:

$$A_{1/2}(\tau) = \frac{2}{\tau^2} (\tau + (\tau - 1) f(\tau)) \quad ; \quad A_1(\tau) = -\frac{1}{\tau^2} (2\tau^2 + 3\tau + 3(2\tau - 1)f(\tau)) \quad (3.7)$$

with:

$$f(\tau) = \begin{cases} \arcsin^2 \sqrt{\tau}, & \tau \leq 1 \\ -\frac{1}{4} \left[\log \frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi \right]^2, & \tau > 1 \end{cases}. \quad (3.8)$$

In the limit $\tau \rightarrow 0$, $A_{1/2} \rightarrow 4/3$ and $A_1 \rightarrow -7$, as expected from the threshold correction of a massive Dirac fermion and vector boson, respectively. In the holomorphic approximation, the contributions A_{Extra}^{gg} and $A_{Extra}^{\gamma\gamma}$ in (3.5) from the extra states are given in terms of an effective beta function coefficient:

$$A_{Extra}^{gg} = 2\tilde{b}_{SU(3)} \cdot \cos(\alpha + \beta); \quad A_{Extra}^{\gamma\gamma} = \tilde{b}_{EM} \cdot \cos(\alpha + \beta) \quad (3.9)$$

where $\tilde{b}_G = b_G \frac{v^2}{\Lambda_G^2}$ with Λ_G a characteristic scale for the states charged under gauge group G . The factor of two in A_{Extra}^{gg} is due to the relative factor of $C_2(fund) = 1/2$ appearing in the $SU(3)$ beta function contribution from the SM states. Thus, γ_{gg} depends on the three parameters $\{\tilde{b}_{SU(3)}, \alpha, \beta\}$, while $\gamma_{\gamma\gamma}$ depends on $\{\tilde{b}_{EM}, \alpha, \beta\}$.

3.2 LHC Constraints

Using our analysis of the contributions of the extra sector to the Higgs partial widths, we now study constraints from the LHC. See also related studies of constraints for various extensions of the Standard Model such as 2HDM models [41–43], fourth generation models [44, 45], and radion models [46, 47]. To frame our discussion, let us first recall the main Higgs search channels which have been studied so far. Both ATLAS and CMS report an excess near 125 GeV coming from gluon fusion production, with subsequent decay to either via

¹⁰In the numerical analysis we also include subleading contributions from SM states to the width of the Higgs.

$h^0 \rightarrow \gamma\gamma$, $gg \rightarrow h^0 \rightarrow ZZ^*$. Additionally, CMS reports a $\gamma\gamma jj$ channel, which will contain contributions from both vector boson fusion and gluon fusion when two extra forward jets are radiated.

While still preliminary, the present limits indicate a signal which is roughly consistent with a Standard Model-like Higgs. Using the notation in (3.3), we use the following experimental values for the channels $pp \rightarrow h^0 \rightarrow \gamma\gamma$ ($R_{g\gamma}$), $pp \rightarrow h^0 \rightarrow ZZ^*$ (R_{gZ}) and $pp \rightarrow h^0 jj \rightarrow \gamma\gamma jj$ which has contributions from both vector boson fusion ($R_{V\gamma}$) and gluon fusion:

$$R_{g\gamma} = 1.4^{+0.7}_{-0.7}, \quad R_{gZ} = 0.8^{+0.8}_{-0.4} \quad (3.10)$$

$$\frac{R_{V\gamma} + \frac{\eta}{2.6} R_{g\gamma}}{1 + \frac{\eta}{2.6}} = 3.7^{+2.5}_{-1.8} \quad (3.11)$$

where for the first two channels, we use the combined results from ATLAS and CMS and for the third we use the CMS result (see e.g. [43] and also [48, 49]). The expression for the $\gamma\gamma jj$ channel (3.11) is obtained from the schematic relation:

$$\sigma(pp \rightarrow h^0 jj \rightarrow \gamma\gamma jj) = (A_v^{jj} \sigma_{VBF} + A_g^{jj} \sigma_{ggF}) \times BR_{h^0 \rightarrow \gamma\gamma} \quad (3.12)$$

where σ_{VBF} and σ_{ggF} are respectively the vector boson fusion and gluon fusion production cross sections, $BR_{h^0 \rightarrow \gamma\gamma}$ denotes the $h^0 \rightarrow \gamma\gamma$ branching ratio, and A_v and A_g are the acceptances for the $\gamma\gamma jj$ channel associated with these two production channels¹¹.

In this work we will take the quoted numerical values and error bars at face value, and ask what regions of parameter space for a given model with a supersymmetric extra sector are consistent with these values. This leads to a non-Bayesian weighting of the various regions of parameter space, but already provides valuable information about the size of the possible contributions from a supersymmetric extra sector (barring contributions from other MSSM states). Though beyond the scope of the present work, once the statistics of the various channels improve, it would be interesting to do a statistical likelihood analysis for such models, weighted by the significance of the various LHC channels.

Since we have the parametric form for the leading-order contributions of the extra sector to Higgs processes, we can study which regions of parameter space are consistent with these numbers. As mentioned earlier, our main assumption is that all other contributions to $h^0 \rightarrow gg$, $\gamma\gamma$ from MSSM superpartners are decoupled. We also assume that branching

¹¹CMS reports that in the SM, one expects 2.01 events from VBF and 0.76 from gluon fusion with the applied cuts [50]. This gives the ratio $\frac{A_g^{jj} \sigma_{ggF}^{SM}}{A_v^{jj} \sigma_{VBF}^{SM}} \sim \frac{\eta}{2.6}$, with η an order one factor to take into account present uncertainties. For specificity, in all plots we take $\eta = 1$, as this corresponds to the value used in [50]. We have also considered values of η up to 2. Though it does not seem to change the results qualitatively, it does add additional (small) regions to $\gamma\gamma jj$.

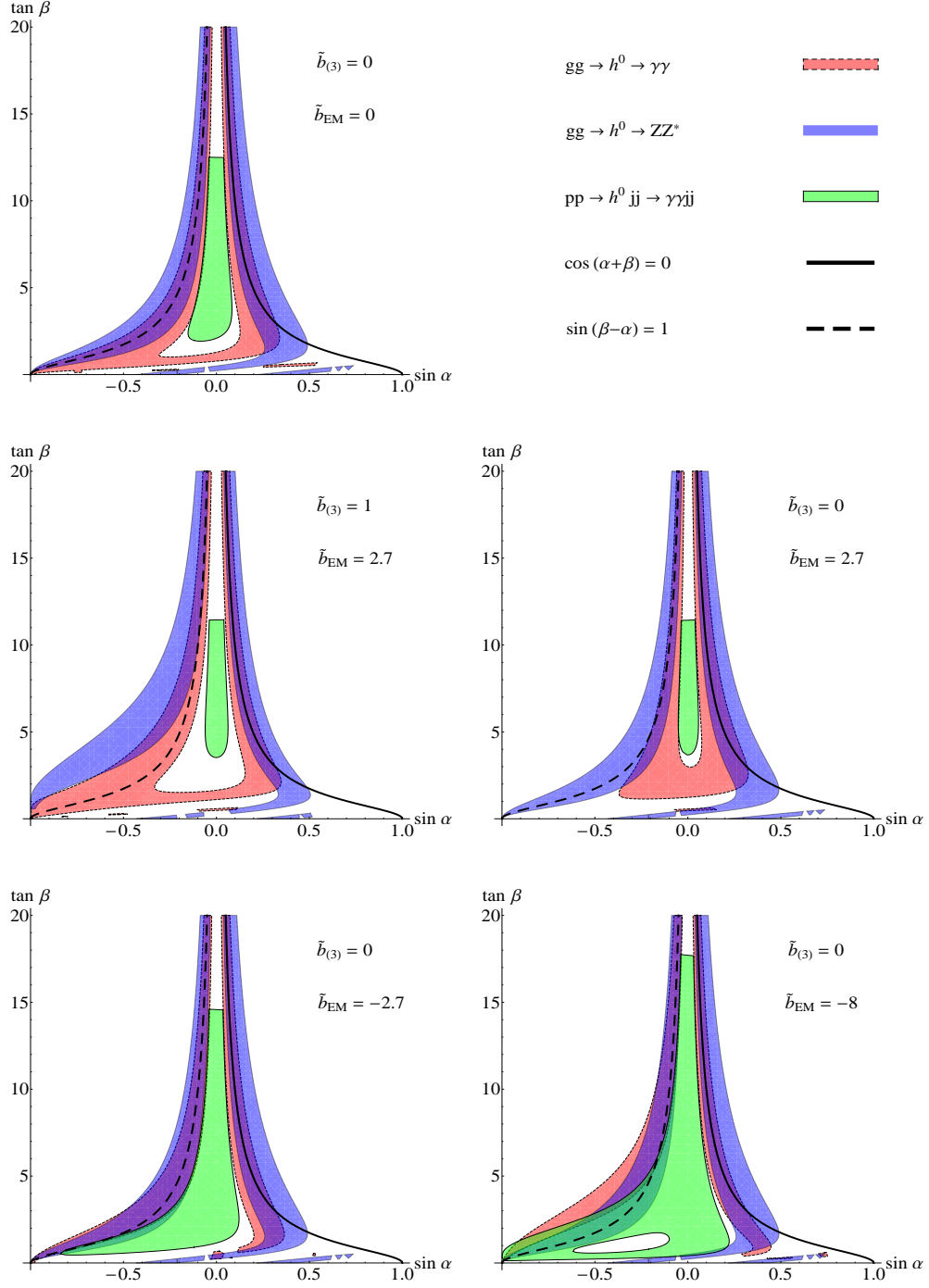


Figure 1: For different values of $\tilde{b}_{SU(3)}$ and \tilde{b}_{EM} , we plot regions in $(\sin \alpha, \tan \beta)$ which are consistent with present limits on the reported LHC signals $gg \rightarrow h^0 \rightarrow \gamma\gamma, ZZ^*$ and $pp \rightarrow h^0 jj \rightarrow \gamma\gamma jj$. See figure 2 for a plot which focuses on the low $\tan \beta$ region.

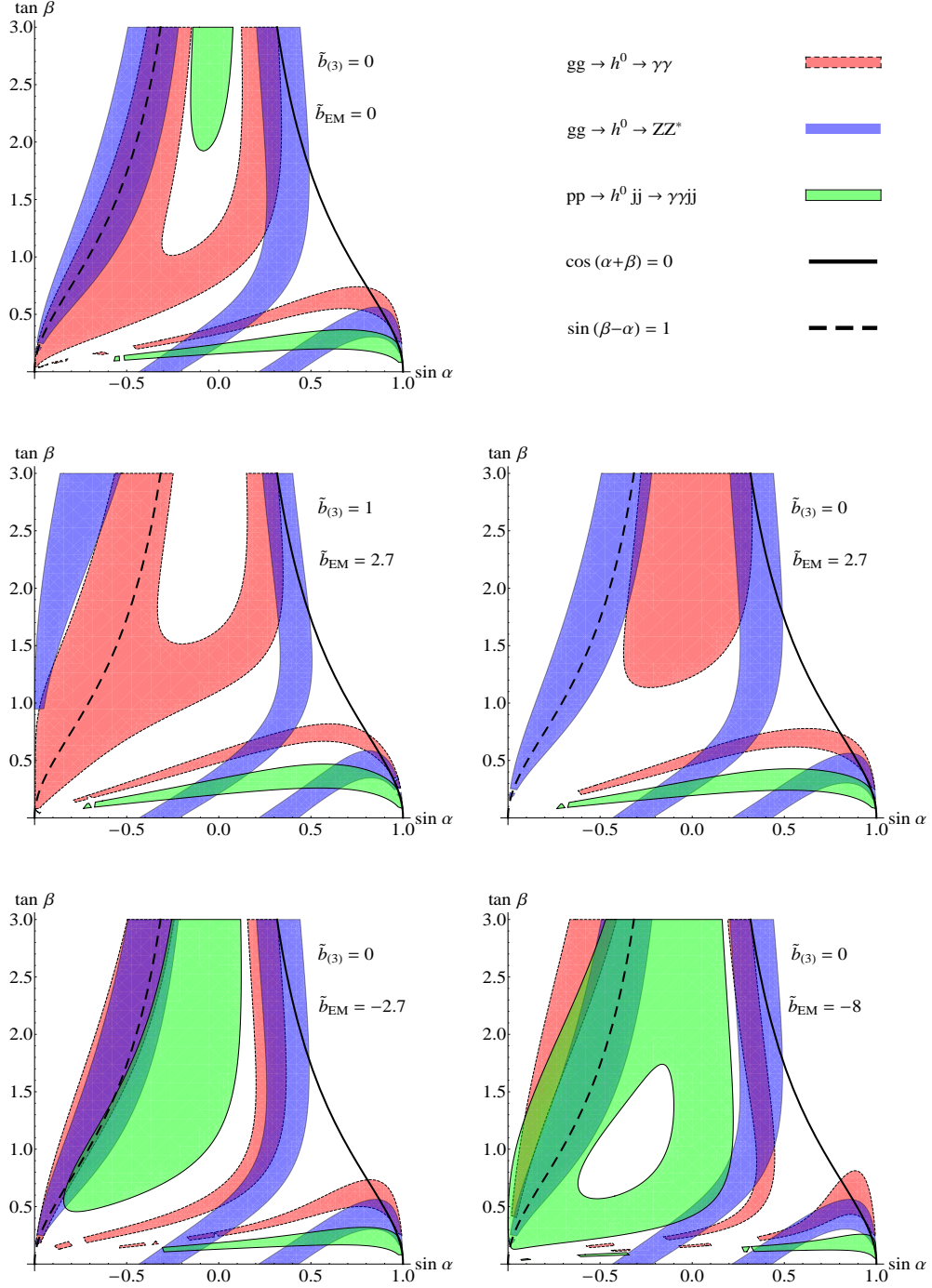


Figure 2: For different values of $\tilde{b}_{SU(3)}$ and \tilde{b}_{EM} , we plot regions in $(\sin \alpha, \tan \beta)$ in the low $\tan \beta$ regime which are consistent with present limits on the reported LHC signals $gg \rightarrow h^0 \rightarrow \gamma\gamma, ZZ^*$ and $pp \rightarrow h^0 jj \rightarrow \gamma\gamma jj$. The case $\tan \beta < 1.2$ is theoretically disfavored, though it is interesting to see that present searches are still consistent with small slivers in this range.

fractions to SM singlets of the extra sector are a subleading contribution. Figure 1 shows the regions in the $(\sin \alpha, \tan \beta)$ plane which are consistent with the experimental values in (3.10) and (3.11) for various values of the effective coefficients $\tilde{b}_{SU(3)}$ and \tilde{b}_{EM} as in equation (3.9). See figure 2 for a plot focusing on the $\tan \beta < 3$ region. For extra sectors which retain gauge coupling unification $b_{SU(3)} = \frac{3}{8} b_{EM}$, but $\tilde{b}_{SU(3)}$ could still be different from $\frac{3}{8} \tilde{b}_{EM}$ if the scales $\Lambda_{SU(3)}$ and Λ_{EM} are different. We find similar behavior when $\tilde{b}_{SU(3)} = 0$ but $\tilde{b}_{EM} \neq 0$, as can happen if the colored states of an extra sector have been decoupled.

To interpret figures 1 and 2, it is helpful to focus on the two limits which exhibit decoupling behavior. The first is the well-known 2HDM decoupling limit $\sin(\beta - \alpha) = 1$, where only h^0 has tree level couplings to the vector bosons. The other limiting case corresponds to $\cos(\alpha + \beta) = 0$, where the extra sector states do not contribute to $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$. By inspection of figure 1, a majority of the parameter space from gluon fusion production is compatible with both limits, but only small slivers are also compatible with $\gamma\gamma jj$. Note that naively one might have thought that it is possible to increase the branching fraction for $h^0 \rightarrow \gamma\gamma$ by lowering the total width of the Higgs through a reduction in $h^0 \rightarrow b\bar{b}$, but much of the parameter space where this could work is already disfavored by current data. Switching on $\tilde{b}_{EM} > 0$ decreases the $h^0 \rightarrow \gamma\gamma$ decay rate because this term destructively interferes with the one arising from the W -loop, whereas the branching fraction appears to be higher than in the Standard Model. However, as shown in figure 2, there are small pockets at $\tan \beta < 1$ which are still viable¹². Note that these regions are close to the curve $\cos(\alpha + \beta) = 0$, implying that the effect of the extra states is suppressed despite a non-negligible $\tilde{b}_{SU(3)}, \tilde{b}_{EM}$. The fact that only small regions are allowed for positive \tilde{b}_G is significant, because as remarked in section 2, unitarity demands $\tilde{b}_G > 0$ when the extra sector is a supersymmetric CFT.

The case $\tilde{b}_G < 0$ is also of interest, though it does not describe a conformal extra sector. This can happen when the Standard Model gauge group embeds in a larger gauge group which contains massive $U(1)_{EM}$ charged vector bosons, as for example in various left-right symmetric extensions of the Standard Model. Here, we observe that it is much easier to remain in accord with present experimental constraints on Higgs searches. This is to be expected, for now the states of the extra sector add constructively with the W -loop in the $h^0 \rightarrow \gamma\gamma$ channel.

Consider next the other modes of the 2HDM sector, H^0 , A^0 , and H^\pm . The corresponding bounds in this case are much more model-dependent. Details of the relative mass spectra, mixing angles, and possible CP-violating contributions will all enter in the analysis of the possible signals of this sector, not to mention the additional contributions from the extra states. Thus, here we confine our discussion to some general comments. If these modes are

¹²Such regions are somewhat problematic from a theoretical standpoint, because they require a large top quark Yukawa, which in turn leads to a Landau pole for this Yukawa at low scales. For such reasons, it is common to impose the condition $\tan \beta > 1$.

heavier than around 250 GeV, then various decay modes such as $A^0 \rightarrow h^0 Z$, $H^0 \rightarrow h^0 Z$, $H^0 \rightarrow h^0 h^0$ could be important for small to moderate $\tan\beta$ [39]. If one is not far from the decoupling limit of the SUSY 2HDM, then the couplings of H^0 to WW or ZZ can be suppressed relative to the SM. These two effects could easily allow one to evade the current bounds from ATLAS and CMS in the WW and ZZ channels, which have been used to put limits on the Higgs cross-sections for such masses [1, 2].

At tree level, the CP-odd state A^0 does not couple to massive vector bosons at all, so there are no bounds for A^0 from these channels. For H^0 , A^0 heavier than about 350 GeV, decays into $t\bar{t}$ will dominate for small $\tan\beta$, so $t\bar{t}$ resonance searches could impose additional limits on $\sigma \cdot BR(t\bar{t})$ [51]. However, the current bounds on $\sigma \cdot BR(t\bar{t})$ for e.g. a 400 GeV resonance decaying into $t\bar{t}$ are quite mild, around 30-40 pb , which is much larger than the MSSM production cross-section of H^0 and A^0 with a mass of 400 GeV. When $\tilde{b}_{SU(3)} > 0$, the production of the heavy states will typically be enhanced relative to a comparable mass h^0 . This is evident for all mixing angles in the case of A^0 , and for H^0 this is the case when $\cos(\alpha + \beta) = 0$, which is the limit where loop contributions to h^0 processes decouple. For moderate values of $\tilde{b}_{SU(3)}$, this enhancement is still not large enough to be an issue, but future data will provide further constraints for such cases with small or moderate $\tan\beta$. For large $\tan\beta$, decays into $b\bar{b}$ are the dominant modes, which are very hard to dig out of the background. It is worth noting that these heavy Higgses could decay into hidden sector singlets providing an invisible decay width, which reduces the branching fraction to visible channels¹³, and further loosens the bounds on such models. Finally, we note that H^\pm mostly tend to decay to $t b$ and $\tau\nu$, and are quite hard to search for. At present, there exist no robust constraints on these states.

To summarize, even though still quite preliminary, the recent hints of an SM-like Higgs already provide an excellent probe into potential signatures of new physics. Within the set of caveats already discussed interesting bounds can be placed on large classes of models, especially for $\tilde{b}_G > 0$, which includes superconformal extra sectors.

4 Examples

In the interest of concreteness, in this section we provide some specific models for extra sectors which can couple to the visible sector. Our aim here is not to construct fully viable phenomenological models, but rather to illustrate that the assumptions necessary to utilize the holomorphic approximation can be met. We also illustrate that there are examples where the Higgs-extra sector Yukawas can be large, but the shift in the Higgs anomalous dimensions are small. This can occur because of cancellations between various contributions to the Higgs anomalous dimensions. As simple cases, we begin with a weakly

¹³Note that relative to h^0 , this is a more natural possibility for the heavy Higgses.

coupled model, and then consider an SQCD-like example. As another class of examples, we discuss some string-inspired SCFTs which evade most of the issues (e.g. inducing low scale Landau poles in the visible sector) which afflict SQCD-like extra sectors. Finally, we note that in non-conformal cases it is possible to have negative b_G .

4.1 Weak Coupling

Let us illustrate the general pattern of Higgs mass dependence in a simple example, with some additional vector-like quark superfields Q, U, D and $\tilde{Q}, \tilde{U}, \tilde{D}$ which couple to the Higgs fields H_u and H_d via:

$$W = \lambda_u H_u Q U + \lambda_d H_d Q D + \tilde{\lambda}_u H_u \tilde{Q} \tilde{D} + \tilde{\lambda}_d H_d \tilde{Q} \tilde{U} \quad (4.1)$$

$$+ \frac{M_Q}{\sqrt{2}} Q \tilde{Q} + \frac{M_U}{\sqrt{2}} U \tilde{U} + \frac{M_D}{\sqrt{2}} D \tilde{D} \quad (4.2)$$

in the obvious notation. Turning on vevs for the Higgs fields, the holomorphic mass matrix splits into up-type and down-type pieces:

$$\mathcal{M}_u = \frac{1}{2\sqrt{2}} \begin{bmatrix} 0 & \lambda_u v_u & M_Q & 0 \\ \lambda_u v_u & 0 & 0 & M_U \\ M_Q & 0 & 0 & \tilde{\lambda}_d v_d \\ 0 & M_U & \tilde{\lambda}_d v_d & 0 \end{bmatrix}, \quad \mathcal{M}_d = \frac{1}{2\sqrt{2}} \begin{bmatrix} 0 & \lambda_d v_d & M_Q & 0 \\ \lambda_d v_d & 0 & 0 & M_D \\ M_Q & 0 & 0 & \tilde{\lambda}_u v_u \\ 0 & M_D & \tilde{\lambda}_u v_u & 0 \end{bmatrix}. \quad (4.3)$$

where our basis of fields for the two matrices is $(U_L, U_R, \tilde{U}_L, \tilde{U}_R)$ and $(D_L, D_R, \tilde{D}_L, \tilde{D}_R)$. The determinant of the two matrices is:

$$\det \mathcal{M}_u = \frac{1}{64} \left(M_Q M_U - \lambda_u \tilde{\lambda}_d v_u v_d \right)^2, \quad \det \mathcal{M}_d = \frac{1}{64} \left(M_Q M_D - \tilde{\lambda}_u \lambda_d v_u v_d \right)^2. \quad (4.4)$$

In this case, all states get a mass which depends on the Higgs vev. This can be seen by working in the limit $M_Q, M_D, M_U \rightarrow 0$. One can also read off the corresponding contribution to the Standard Model beta functions; the thresholds are $b_{SU(3)} = 4$, $b_{SU(2)} = 3$ and $b_{U(1)} = 11/5$, where we have adopted an $SU(5)_{GUT}$ normalization of $U(1)_Y$. Hence, $b_{EM} = b_{SU(2)} + \frac{5}{3} b_{U(1)} \sim 6.66$. In order to achieve an effective $\tilde{b}_{EM} \sim 2.7$ one requires a characteristic scale $\Lambda_{EM} \sim 2v \sim 490$ GeV. Of course, the mass of the extra states depends on the sizes of the Yukawas, a feature we have absorbed into our convention for Λ_{EM} . It should be clear that this example can easily be extended to include full GUT multiplets. Finally, let us note that for a weakly coupled model such as this one, incorporating supersymmetry breaking effects can shift the relative masses of the scalars and fermions. This is because there is no conformal suppression of such soft breaking terms here. Note, however, that even if the scalars get large soft supersymmetry breaking masses, a remnant of the holomorphic

approximation persists in the form of the contribution from the fermions to $h^0 \text{Tr}_G F^2$.

4.2 An SQCD-Like Model

We now move on to an example where the extra sector is an SQCD-like theory. We study the anomalous dimensions of the various fields with and without the Higgs sector couplings, and the consequent change these anomalous dimensions induce in the visible sector beta functions.

Consider an extra sector with gauge group $SU(N_c)$ and matter fields $L_u^{(i)} \oplus L_d^{(i)}$ in the $(2_{-1/2}, N_c) \oplus (2_{1/2}, \overline{N_c})$ of $SU(2)_L \times U(1)_Y \times SU(N_c)$, where the flavor index $i = 1, \dots, N_f$. We also introduce a pair of singlets $S_u \oplus S_d$ in the $(1_0, \overline{N_c}) \oplus (1_0, N_c)$, so that we can have nontrivial interactions between the extra sector and the Higgs. Note that since the states of the extra sector are only charged under $SU(2)_L \times U(1)_Y$, the resulting threshold corrections will not affect the leading-order gluon fusion cross section, but will alter the $h^0 \rightarrow \gamma\gamma$ decay channel.

Without a superpotential, this theory is just $SU(N_c)$ SQCD with $2N_f + 1$ flavors and H_u and H_d are decoupled free fields. Here we are interested in a conformal extra sector so we take $\frac{3}{2}N_c < 2N_f + 1 \leq 3N_c$, to remain in the conformal window of SQCD. The resulting R-charges are

$$R_H = \frac{2}{3}, \quad R_S = R_L = 1 - \frac{N_c}{2N_f + 1}. \quad (4.5)$$

The dimension Δ of the scalar component of a chiral primary superfield is related to the R-charge via the formula $\Delta = 3R/2$. Adopting an $SU(5)_{GUT}$ normalization of the $U(1)_Y$ generator, the threshold correction from the extra states to the $SU(2)_L$ and $U(1)_Y$ beta functions is:

$$b_{SU(2)} = -3 \times N_c N_f (R_L - 1) = \frac{3N_c^2 N_f}{2N_f + 1}, \quad (4.6)$$

$$b_{U(1)} = -3 \times \frac{3}{10} N_c N_f (R_L - 1) = \frac{9N_c^2 N_f}{20N_f + 10} \quad (4.7)$$

while the contribution to $b_{EM} = b_{SU(2)} + \frac{5}{3}b_{U(1)}$ is:

$$b_{EM} = -\frac{9}{2} N_c N_f (R_L - 1) = \frac{9}{2} \times \frac{N_c^2 N_f}{2N_f + 1}. \quad (4.8)$$

Now consider switching on the superpotential interaction

$$W_{mix} = \lambda_i H_u L_u^{(i)} S_u + \tilde{\lambda}_j H_d L_d^{(j)} S_d \quad (4.9)$$

which can induce a flow to a new interacting fixed point. As we will shortly verify, these mixing terms can be large but nevertheless produce only a small shift in the scaling dimensions of the Higgs fields. Let us now proceed to an analysis of the IR fixed point in the presence of W_{mix} .

As can be checked, there are still only three independent R-charges R_H , R_S and R_L . Along with the condition that the R-symmetry be anomaly-free, enforcing that the superpotential be marginal gives two conditions on three undetermined R-charges. Maximizing $a = \frac{3}{32} (3\text{Tr}R^3 - \text{Tr}R)$ [28] over the remaining variable, we obtain the R-charge assignments:

$$R_H = \frac{y+x}{z}, R_S = 1 + \frac{N_c - 2N_f R_H}{2N_f - 1}, R_L = 1 - \frac{N_c - R_H}{2N_f - 1} \quad (4.10)$$

where:

$$x = \sqrt{9 \left(1 - (4 + N_c^2)N_f + 4N_f^2\right)^2 + 8(2N_f - 1)^2(-1 + N_f(4 + N_c + 2N_f(N_c - 2)))} \quad (4.11)$$

$$y = -3 + 3N_f(4 + N_c^2) - 12N_f^2 \quad (4.12)$$

$$z = -3 + 3N_f(4 + N_c + 2N_f(N_c - 2)). \quad (4.13)$$

With our modified R-charge assignments, we can recompute the values of the scaling dimensions, and the changes to the beta functions. In this case, it is important to include the fact that the dimension of the Higgs will now be shifted away from its free field value. The contribution of the extra sector states to the $SU(2)_L$ and $U(1)_Y$ beta functions will in this case be:

$$b_{SU(2)} = -3 \times N_c N_f (R_L - 1) - 3 \times (R_H - 1) + 3 \times \left(\frac{2}{3} - 1\right) \quad (4.14)$$

$$b_{U(1)} = -3 \times \frac{3}{10} N_c N_f (R_L - 1) - 3 \times \frac{3}{10} (R_H - 1) + 3 \times \frac{3}{10} \times \left(\frac{2}{3} - 1\right) \quad (4.15)$$

where in the above, we have also included the contribution from a shift in the dimension of the Higgs away from its free field value.¹⁴ Finally, the contribution to b_{EM} is:

$$b_{EM} = -\frac{9}{2} N_c N_f (R_L - 1) - \frac{9}{2} \times \left(R_H - \frac{2}{3}\right). \quad (4.16)$$

As an example, we can take $N_c = 2$, and $N_f = 2$, which yields $\Delta_H = 1.15$, $\Delta_S = 0.97$, $\Delta_L = 0.88$ and $b_{EM} = 6.9$. Comparing the value of b_{EM} without mixing to the case with mixing, we see that $\delta b_{EM}/b_{EM} \sim 0.03$, which justifies the use of the holomorphic approximation.

¹⁴Alternatively, one can simply consider the full $H \oplus L$ contribution in both the UV and the IR. Note that the difference between the UV and IR contributions will be the same, however.

Switching on vector-like mass terms to decouple the extra sector, an effective $\tilde{b}_{EM} \sim 2.7$ requires $\Lambda \sim 400$ GeV.

It is also of interest to study Banks-Zaks fixed points to find additional regimes where the Higgs dimension only shifts by a small amount. For example, taking $N_f = \frac{3}{2}(N_c - 1)$ (which is just below the asymptotic freedom bound $2N_f + 1 = 3N_c$) and expanding in the large N_c limit, we have:

$$\Delta_H = 1 + \frac{2}{3} \frac{1}{N_c} + O\left(\frac{1}{N_c^2}\right) \quad (4.17)$$

and $\delta b_{EM}/b_{EM} \sim 1.8/N_c^4$. Of course, in this case, there is also a significant increase in the beta functions; we have $b_{EM} \sim 2.25 \times N_c^2$ in the mixed theory which will lead to a low-scale visible sector Landau pole.

An unappealing feature of this example is that the matter fields do not form GUT multiplets, so there is no chance for gauge coupling unification. Similar issues confront large rank SQCD-like extra sectors, because they lead to low scale Landau poles. This leaves only a few low rank gauge groups. This, and other issues can be overcome in recently studied CFTs arising in explicit string constructions [25].

4.3 String-Inspired Example

We now turn to some examples based on a strongly coupled limit of IIB string theory known as “F-theory”. From a field theory standpoint, these F-theory CFTs can be viewed as $\mathcal{N} = 1$ deformations of an $\mathcal{N} = 2$ SCFT with an E_8 flavor symmetry, related to the famous Minahan-Nemeschansky SCFTs [52, 53]. The lowest dimension chiral primaries of the $\mathcal{N} = 2$ Minahan-Nemeschansky theory are a dimension two operator \mathcal{O}_{248} in the adjoint of E_8 , and a dimension six operator Z which is a singlet under E_8 . The \mathcal{O}_{248} ’s are the analogue of mesons in SQCD-like theories. When Z gets a vev, all of the charged states pick up a vector-like mass.

Relevant $\mathcal{N} = 1$ deformations of the $\mathcal{N} = 2$ Minahan Nemeschansky theory lead to $\mathcal{N} = 1$ theories, where the mass deformations transform in the adjoint of E_8 . Such deformations initiate a breaking pattern down to $SU(5)_{GUT}$. Promoting the remaining mass deformations to Standard Model fields yields couplings such as $H_u \mathcal{O}_u$ and $H_d \mathcal{O}_d$. These deformations correspond to marginal couplings in the infrared, and can in principle be large. However, the contribution to the Higgs anomalous dimension can still be small [25]. See [10, 15, 25, 54–56] for studies of formal and phenomenological aspects of such extra sectors.

These theories automatically overcome many of the issues which one typically encounters in SQCD-like theories. For example, the resulting contribution to the visible sector beta functions tends to be much smaller than in SQCD-like theories. This is basically because the dynamics of the theory is governed (on the Coulomb branch) by a strongly coupled

$U(1)$ gauge theory, rather than by a non-abelian gauge theory with high rank. Once the Higgs gets a vev, all of the states charged under $SU(5)_{GUT}$ get a mass proportional to the Higgs vevs.¹⁵ As a consequence, the Higgs would be expected to have decays to visible sector states, with negligible invisible width.

The analysis of operator scaling dimensions and the value of the beta functions has been studied in detail in [25, 56], so we shall simply summarize some examples in what follows. Consider first a “ S_3 monodromy scenario”. The values of the beta function coefficients b_G without Higgs-CFT mixing (b^{UV}), and with Higgs-CFT mixing (b^{MIX}) are:

$$b^{UV} = \frac{3k_{E_8}}{4}t^{UV}, \quad b^{MIX} = \frac{3k_{E_8}}{4}t^{MIX} \quad (4.18)$$

where here we have assumed no additional mixing between the visible sector and the D3-brane, and we have dropped the subleading GUT distorting contributions to the beta functions. The parameter $k_{E_8} = 12$ in the $\mathcal{N} = 2$ E_8 Minahan-Nemeschansky theory (see e.g. [57]). In this case, $t^{UV} \sim 0.40$ and $t^{MIX} \sim 0.36$, as found in [25] and [56], respectively. In the absence of electroweak symmetry breaking, H_u and H_d become operators of the IR SCFT, with scaling dimensions $\Delta_{H_u} = \Delta_{H_d} = 1.08$, which indicates low Higgs-extra sector mixing. The value of $\delta b/b$ (for an $SU(5)_{GUT}$ beta function) is $\delta b/b \sim 0.1$, which justifies our approximation. Note that $b_{GUT} \sim 3.2$ and $b_{EM} \sim 8.5$. Introducing a vector-like mass for the states by going onto the Coulomb branch of the theory, achieving $\tilde{b}_{SU(3)} \sim 1$ and $\tilde{b}_{EM} \sim 2.7$ requires a characteristic scale of order $\Lambda_{SU(3)} \sim \Lambda_{EM} \sim 440$ GeV.

As another class of examples, we can consider the “ $Dih_4^{(2)}$ monodromy scenario”. The values of the parameters in this case are $t^{UV} \sim 0.29$ and [25] $t^{MIX} \sim 0.27$ [56]. In this case, the coupling $H_u \mathcal{O}_u$ is actually irrelevant, and H_u remains of dimension one, while H_d has dimension $\Delta_{H_d} = 1.02$. The value of $\delta b/b \sim 0.07$, and the overall value of the beta function coefficients are $b_{GUT} \sim 2.4$ and $b_{EM} \sim 6.4$. In this case, an effective $\tilde{b}_{SU(3)} \sim 1$ and $\tilde{b}_{EM} \sim 2.7$ requires $\Lambda_{SU(3)} \sim \Lambda_{EM} \sim 380$ GeV.

4.4 Non-Conformal Theories

In any conformal theory, the sign of b_G is constrained by unitarity to be positive, as is true in the above examples. However, if the extra sector is in a non-conformal phase, b_G may

¹⁵One can see this in a variety of ways. Geometrically, these SCFTs are realized by a D3-brane probing an E-type point of the F-theory geometry. The Standard Model chiral superfields correspond to modes localized at the intersection of two intersecting seven-branes, and the D3-brane sits at the intersection of several such branes. When the Higgs gets a vev, the branes recombine, and move away from the location of the D3-brane. This gives a mass to the $SU(5)_{GUT}$ charged states, i.e. the “3-7_{vis}” strings, as well as all “3-7_{hid}” strings. In more field theoretic terminology, giving the Higgs a vev allows one to form a mass deformation of the original $\mathcal{N} = 2$ theory with a non-trivial Casimir invariant under E_8 . This in particular means that all states charged under the original E_8 flavor symmetry will pick up a mass.

now take either sign. It is straightforward to see how one could get a beta function with opposite sign: Since gauge bosons contribute negatively to b_G while matter contributes positively, one just needs a regime in which the contribution from the vectors outweighs the contribution from the scalars. One example is the left-right symmetric model of [58, 59], in which there is an extra $SU(2)_R$ gauge group which gets Higgsed. A W' running in the loop will contribute with the same sign as a W , which can tend to enhance $h^0 \rightarrow \gamma\gamma$. Of course, to get enough of an enhancement may require multiple W' s, since there is a generic suppression of order v^2/Λ^2 for Λ on the order of the mass of the W' . Further, one can expect additional constraints from other considerations. We leave it as an open problem in model building whether a sufficiently large enhancement to $h^0 \rightarrow \gamma\gamma$ with W' 's can be achieved.

Actually, the case of left-right symmetry breaking is instructive for a more theoretical reason, because it would seem to violate the mixing angle dependence we argued should hold in the holomorphic approximation. Indeed, the $h^0 W'^+ W'^-$ vertex is proportional to $\sin(\beta - \alpha)$, which is certainly different from $\cos(\alpha + \beta)$. Note, however, that to remain in the holomorphic approximation, one must satisfy the D-term equation of motion for $SU(2)_R$. In the limit where the Higgs fields are the sole source of mass for the extra sector, we have the D-flatness condition $v_u = v_d$ so that $\beta = \pi/4$. In this special case, we have $\cos(\alpha + \beta)|_{\beta=\pi/4} = \sin(\beta - \alpha)|_{\beta=\pi/4} = (\cos \alpha - \sin \alpha)/\sqrt{2}$. With additional sources of mass terms, there are more general, model dependent ways to satisfy the D-term constraints which would be interesting to consider as well.

5 Conclusions

In this paper we have developed a general method for extracting the contributions to the processes $h^0 \rightarrow gg$ and $h^0 \rightarrow \gamma\gamma$ from an approximately supersymmetric extra sector which mixes with the Higgs. We have explained how holomorphy constrains the dimension five operators, and in particular, fixes the dependence on the Higgs mixing angles. Further, when the Higgs is the sole source of mass for a superconformal sector, we have seen that the effects of the extra sector are fully specified by *calculable* coefficients. Applying these observations, we have explained how to calculate the contribution to various Higgs processes from such scenarios, how LHC data provides constraints on the properties of such extra sectors, and moreover, have given explicit examples where the assumptions of the holomorphic approximation can be met. In the remainder of this section, we discuss some potential avenues of future investigation.

From a phenomenological point of view, it would be very interesting to study the signatures of the colored and electroweak states in detail, so that one can devise carefully designed searches to look for them. This is especially true if future Higgs measurements,

when interpreted within this framework, suggest a value of Λ_G in an experimentally accessible range.

From a theoretical point of view, it would be interesting to see whether additional calculable information about extra sectors could also be extracted and repackaged in terms of higher dimension operators involving Higgs fields.

An important technical assumption of this work has been that the extra sector is approximately supersymmetric. While less quantitative control is available in the non-supersymmetric case, it is also clearly more general. Phrased differently, one can view our computation as a guide for “how far” from supersymmetric an extra sector must deviate in order to evade the parametric form found here. Characterizing the form of possible (small) deviations from the holomorphic approximation would clearly be of interest. Related to this, it would be of interest to consider a more general phenomenological analysis of the Higgs away from the vector-like mass limit of the extra sector.

Finally, anticipating the significant improvement in experimental Higgs search channels expected by the end of even 2012, it would of course be interesting to later return to a more detailed study of potential constraints and evidence for the parametric form of couplings expected in the holomorphic Higgs regime.

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